

# PRIMORDIAL BLACK HOLES FROM NON-EQUILIBRIUM SECOND ORDER PHASE TRANSITION

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## Abstract

The collapse of sufficiently large closed domain wall produced during second order phase transition in the vacuum state of a scalar field can lead to the formation of black hole. The origin of domain walls with appropriate size and energy density could be a result of evolution of an effectively massless scalar field at the inflational epoch. We demonstrate that in this case the situation is valued when there are compact domains of less favorable vacuum surrounded by a sea of another vacuum. Each domain has a surface composed of vacuum wall that stores a significant amount of energy, and can collapse into the black hole. This offers the way of massive primordial black holes formation in the early Universe.

## 1 Introduction

It is well known that any object put within its gravitational radius forms a black hole (BH). At present time BHs can be naturally created only in the result of gravitational collapse of stars with the mass exceeding three Solar masses <sup>1)</sup> in the end of their evolution. From the other hand, it has long been known <sup>2, 3, 4, 5, 6, 7)</sup> that primordial black hole (PBH) formation is possible in the early Universe. PBH can form when the density fluctuations become larger than unity on a scale intermediate between Jeans length and horizon size. Other possibilities are related to the dynamics of various topological defects such as the collapse of cosmic strings <sup>3)</sup> from the thermal second order phase transition or to the collisions of the bubble walls <sup>4, 5)</sup> created at the first order phase transitions. Formally, there is no limit on the mass of PBH that forms after the collapse of highly overdense region, it is only needed to form appropriate spectrum of initial density fluctuations at the inflation <sup>8)</sup>. We cannot expect *a priori* the same "no-mass-limit" condition in the case of PBHs formed by topological defects, because the mass of such PBHs is defined by the correlation length of the respective phase transition.

In this paper we concern with the possibility to form PBHs after the self-collapse of closed domain walls created during a second order phase transition. Such PBHs should have small masses in the case of thermal second order phase transition with usual equilibrium initial conditions. It takes place because the characteristic size of walls coincides with the correlation length of phase transition, which initiate the formation of that domain walls. Usually we deal with high temperature phase transition in the early Universe, that makes the respective correlation length quite small and does not allow to store a significant amount of energy within the individual closed vacuum wall. Moreover the

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phase transition can have quite complicated dynamics, namely it can be accompanied by another phase transition that creates another type of topological defects, namely, vacuum strings<sup>9, 10)</sup> (see discussion in Section 2). As a result we get the mesh of vacuum strings connected by walls where the probability of the existence of closed walls is strongly suppressed<sup>10)</sup>. We discuss here the non-equilibrium scenario of closed vacuum walls formation that opens a new possibility for the massive PBHs production in the early Universe.

The starting point is rather general and evident. If a potential of a system possesses at least two different vacuum states there are two possibilities to populate that states in the early Universe. The first one is that the Universe contains both states populated with equal probability, that takes place under the usual circumstances of thermal phase transition. The other possibility corresponds to the case when two vacuum states are populated with different probability and there are islands of the less probable vacuum, surrounded by the sea of another, more preferable, vacuum. The last possibility can take place when we go beyond the equilibrium conditions, established by pure thermal dynamics. More definitely, it is necessary to redefine effectively the correlation length of the scalar field that drives a phase transition and consequently the formation of topological defects. We will show that the only necessary ingredient for it is the existence of an effectively flat direction(s), along which the scalar potential vanishes during inflation. Then the background de-Sitter fluctuations of such effectively massless scalar field could provide non-equilibrium redefinition of correlation length and give rise to the islands of one vacuum in the sea of another one. In spite of such redefinition the phase transition itself takes place deeply in the Friedman–Robertson-Walker (FRW) epoch. After the phase transition two vacua are separated by a wall, and such a wall can be very big. The motion of closed vacuum walls has been first driven analytically in<sup>11)</sup>. At some moment after crossing horizon they start shrinking due to surface tension. As a result, if the wall does not release the significant fraction of its energy in the form of outward scalar waves, almost the whole energy of such closed wall can be concentrated in a small volume within its gravitational radius what is the necessary condition for PBH formation.

The mass spectrum of the PBHs which can be created by such a way depends on the scalar field potential which parametrizes the flat direction during inflation and triggers the phase transition at the FRW stage. Through the paper we will deal with so called pseudo Nambu–Goldstone (PNG) potential, that is quite common for the particle physics models.

The plan of this paper is as following. In Section 2 the origin of vacuum walls at the PNG potential is discussed. The possibility to form the non-equilibrium conditions for the second order phase transition at the inflantional stage is considered in Section 3. In Section 4 the minimal conditions of PBHs formation from collapsing closed vacuum walls are discussed. The mass spectrum of PBHs is evaluated numerically for the certain choises of parameters.

## 2 Domain Walls at the PNG Potential Under Thermal Approach

Domain walls are planar defects in the vacuum alignment over space. They can appear if the manifold of degenerate (or nearly degenerate) vacua of the theory is disconnected. This is generally the case if there is a discrete symmetry in the potential. Such type of symmetry can be fundamental as, for example, in the double well potential or can arise dynamically during the breakdown induced by instanton effects of a continuous symmetry (see for example<sup>9)</sup>). Moreover, discrete symmetry can be got as a result of explicit breaking of a continuous symmetry due to Yukawa-type couplings as in the case of classical PNG boson or due to higher-order nonrenormalizable interactions that often takes place in string-inspired models<sup>12)</sup>. In all this cases the resulting term of potential has the form

$$V = \Lambda^4 \left( 1 - \cos \frac{\phi}{f} \right) \quad (1)$$

where  $f$  is the scale of  $U(1)$  symmetry breaking. The potential describing a spontaneously broken  $U(1)$  symmetry has standard Mexican hat form with bottom radius  $f$  and reads

$$V(\varphi) = \lambda(|\varphi|^2 - f^2/2)^2, \quad (2)$$

where  $\varphi = \frac{f}{\sqrt{2}} \exp(i\phi/f)$ . The term (1) breaks the continuous symmetry down to the symmetry  $\phi \rightarrow \phi + 2\pi n f$ , where  $n$  is integer. Thus potential (1) has a number of discrete degenerate minima. That minima also can be non-equivalent if the manifold spanned by  $\phi$  is non-compact. Moreover, it is not even necessary for the non-equivalent minima to be exactly degenerate, so the symmetries we consider may be only approximate ones, although non-degeneracy may change the details of wall evolution<sup>12)</sup>. However, as we have noted, many reasonable models can give rise to cosine potentials like (1), therefore we are going to use it in the explicit calculations in this paper. The equation of motion that corresponds to (1) admits kink-like, domain wall solution, which interpolates between two adjacent vacua. If we will consider that wall as lying in the perpendicular to  $x$  axis then the static domain wall solution between two vacua  $\phi_1 = 0$  and  $\phi_2 = 2\pi f$  is given by

$$\phi_{wall}(x, x_0) = 4f \arctan \left( \exp \frac{x - x_0}{d} \right) \quad (3)$$

where  $x_0$  is the location point of the center of the wall,  $d$  is its width,

$$d = \frac{f}{\Lambda^2} = m^{-1} \quad (4)$$

and  $m$  is the mass of PNG field. The surface energy density of the wall reads<sup>13, 10)</sup>

$$\sigma = \int_{-\infty}^{+\infty} \Lambda^4 \left( 1 - \cos \frac{\phi_{wall}}{f} \right) dx \approx 8f\Lambda^2 \quad (5)$$

The thermal dynamics of a system invoking two successive second order phase transitions, with spontaneous  $U(1)$  symmetry breaking and with the explicit one, results in the formation of a quite sophisticated system of topological defects, network of walls bounded by strings (see for review<sup>10)</sup>). We suppose that  $f \gg \Lambda$ . This means that first the spontaneous  $U(1)$  symmetry breaking takes place in the early Universe at the temperature  $T \simeq f$ . After that moment the Universe is filled with global  $U(1)$  strings<sup>14, 10)</sup> and the phase  $\theta = \phi/f$  of complex scalar field  $\phi$  obeys the standard equation of motion

$$\ddot{\theta} + 3H\dot{\theta} + \frac{dV}{d\theta} = 0 \quad (6)$$

As long as the Hubble expansion is fast enough to provide the domination of friction term over the gradient one, the explicit symmetry is restored. It takes place until the moment when the temperature falls down below the value of  $\Lambda$ . At that period the dynamics of global string network implies that there is approximately only one string per Hubble horizon<sup>15, 14, 10)</sup>, that defines the average distance between strings. The correlation length for the second phase transition with explicit  $U(1)$  symmetry breaking is just equal to the average distance between global strings at the moment, when it takes place and triggers the formation of domain walls<sup>10, 14)</sup>. Thus the typical size of domain walls created during last phase transition corresponds to the size of Hubble horizon at the temperature  $T \simeq \Lambda$ . As it has been pointed out in the Introduction, the closed domain walls could collapse to the black hole by radiating away their asphericity due to the surface tension forces<sup>11, 16)</sup>. If we concern the fate of closed domain walls that could be formed after the last phase transition, we have to compare the gravitational radius  $r_g = 2E/m_p^2$  of the energy  $E$  stored in such a closed shell, with the minimal radius that can be achieved in the result of such collapse. The typical size of a closed wall is given by the correlation length of the last phase transition and reads  $L_c = 0.3m_p/(\sqrt{g_*}\Lambda^2)$  ( $g_*$  counts the

effectively number degrees of freedom in the plasma and according to the standard model can not be much less than  $10^2$ ). Combining it with (5) the gravitational radius will read  $r_g \simeq 10^{-3} \frac{f}{\Lambda^2}$ . From the other hand the numerical simulation of the collapse of closed domain wall at the PNG potential <sup>17)</sup> shows that the minimal radius of shrinking cannot be much smaller than the wall's width (4). This means that it is impossible to concentrate the energy of closed walls under the gravitational radius by self-collapse if we consider the domain walls as the result of usual thermal phase transition. Moreover the system of topological defects contains preferably walls bounded by strings and open walls, while the probability of closed walls formation is strongly suppressed <sup>10)</sup>.

To escape such no-go results of equilibrium dynamics we need to enlarge essentially the distance between strings and to increase the probability of closed walls formation. The first problem can be solved by the inflational blow up of the typical distance of string's separation. As well we show below that the non-equilibrium conditions imposed by background inflation allow us to bias the preference of the theory to choose one vacuum state over another that increase essentially the probability of very large closed walls formation.

### 3 Phase Transition With Non-Equilibrium Initial Conditions

We consider the Universe that, due to the existence of an inflaton, goes through a period of inflation and then settles down to the standard FRW geometry. In addition we introduce a complex scalar field  $\varphi$ , not the inflaton, with a large radial mass  $\sqrt{\lambda}f > H_i$ , that has got Mexican hat potential (2). It makes sure that  $U(1)$  symmetry is already broken spontaneously at the beginning of inflation in our cosmological horizon. Thereby we deal only with the phase of that complex field  $\theta = \phi/f$ , which parametrizes potential (1). Under this condition we come to the conclusion, that the correlation length of second order phase transition with spontaneously broken  $U(1)$  symmetry exceeds the present cosmological horizon, and all global  $U(1)$  strings are beyond our horizon <sup>2</sup>. Furthermore, we assume that

$$m \ll H_i, \quad (7)$$

where  $H_i$  is the Hubble constant during inflation. This condition implies that during inflation the potential energy of field  $\phi$  is much smaller than the cosmological friction term what justifies neglecting the potential until the Universe goes deeply into the FRW phase.

The dynamics of such effectively massless scalar field  $\theta$  existing on de Sitter space breaks into two parts (see for review <sup>18)</sup>). First, there is a classical field  $\theta_0$ , which satisfies the standard classical equation of motion (6). During inflation, and long afterward,  $H_i$  is very large (by assumption) compared to the potential (1). It follows that we can drop the gradient term in the equation of motion (6) and resulting equation is solved by  $\theta_0 = \theta_{N_{max}}$ , where  $\theta_{N_{max}}$  is an arbitrary constant. We emphasize that there is no information in the theory that can fix the value of  $\theta_{N_{max}}$ . It is completely arbitrary. We make standard assumption that our present horizon that has been nucleated at the  $N_{max}$  e-folds before the end of inflationary epoch is embedded in an enormous inflation horizon, created by exponential blow up of a single causal horizon. It follows that  $\theta_{N_{max}}$  will be the same over the inter inflationary horizon. We put  $\theta_{N_{max}} < \pi$  without loss of generality. Second we must consider the quantum fluctuations of the phase  $\theta$  at the de Sitter background. It is well known <sup>19)</sup>, that there are quantum fluctuations produced on the vacuum state of  $\theta$  due to the boundary conditions of de Sitter space. These fluctuations are sometimes referred to as contribution to the "Hawking temperature" of de Sitter space but, in fact, there are not true thermal effects <sup>18)</sup>. It makes the dynamics of phase  $\theta$

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<sup>2</sup>Also we can consider, a complex scalar field with a large initial radial positive mass that during inflation undergo a phase transition to a spontaneously broken  $U(1)$  symmetry <sup>8, 18)</sup>. In such a case the correlation length can be smaller than our cosmological horizon, but significantly larger than in the case of usual thermal dynamics

strongly non-equilibrium leading to the non-thermal distribution of scales populated with different vacuums in the postinflationary Universe. Such fluctuations can be described by "quasiclassical" scalar field, which contains components with wavelengths ranging from the size of particle horizon during inflation  $H_i^{-1}$ , all the way out to the inflation horizon  $H_i^{-1}e^{N_{max}}$ . Thus, when the wavelength of a particular fluctuation becomes greater than  $H_i^{-1}$ , the average amplitude of this fluctuation freezes out at some value due to the large friction term in the equation of motion, whereas its wavelength grows exponentially. In the other words such a frozen fluctuation is equivalent to classical field, discussed earlier, that does not vanish after averaging over macroscopic space intervals. The vacuum contains fluctuations of every wavelengths and hence inflation leads to the creation of new regions containing the classical field of different amplitudes with scale greater than  $H_i^{-1}$ . The average amplitude of such fluctuations for massless field generated during each time interval  $H_i^{-1}$  is <sup>19)</sup>

$$\delta\theta = \frac{H_i}{2\pi f} \quad (8)$$

In the other words the phase  $\theta$  makes quantum step  $\frac{H_i}{2\pi f}$  during every e-fold and in every space volume with characteristic size of the order of the  $H_i^{-1}$ . The total number of steps during time interval  $\Delta t$  is given by  $N = H_i \Delta t$ . Such a motion looks like the one-dimensional Brownian motion. The  $\theta_{N_{max}}$  plays the role of the starting point for this Brownian motion. Thus, the initial domain containing phase  $\theta_{N_{max}}$  increases its volume in  $e^3$  times after one e-fold by definition and hence contains  $e^3$  separate, causally disconnected domains of size  $H_i^{-1}$ . Each domain is characterized by average phase value

$$\theta_{N_{max}-1} = \theta_{N_{max}} \pm \delta\theta. \quad (9)$$

In half of these domains the phases evolve toward  $\pi$  while in the other domains they move toward zero. This process is duplicated in each volume of size  $H^{-1}$  during next e-fold. Now at any given scale  $l = k^{-1}$  the size distribution of the phase value  $\theta$  can be described by Gauss' law <sup>20)</sup>.

$$P(\theta_l, l) = \frac{1}{\sqrt{2\pi}\sigma_l} \exp \left\{ -\frac{(\theta_{N_{max}} - \theta_l)^2}{2\sigma_l^2} \right\}, \quad (10)$$

where the dispersion could be expressed in the following manner

$$\sigma_l^2 = \frac{H^2}{4\pi^2} \int_{k_{min}}^k d \ln k = \frac{H^2}{4\pi^2} \ln \frac{l_{max}}{l} = \frac{H^2}{4\pi^2 f^2} (N_{max} - N_l), \quad (11)$$

As long as condition (7) is satisfied the vacuum is well defined by the distribution (10) of fluctuations around chosen constant  $\theta_{N_{max}}$ . Eventually, long after the end of inflationary epoch,  $H(t)$  decreases so significantly that the gradient term in the equation of motion (6) begins to dominate over the friction term, and  $\theta$  starts to oscillate around the one of degenerated minimum. We refer the time  $t_c$  moment at which the condition (7) is not valid anymore as the time moment when the phase transition is triggered. The key point of our consideration is based on the following observation: At  $t_c$  the field configuration of phase  $\theta$  within the volume of contemporary horizon is uniform and defined by  $\theta_{N_{max}}$ . There are, however, fluctuations due to the quasi-classical random field (10) inside this volume and the spectrum of this fluctuations contains all wavelengths up to biggest cosmological scale. Thus the field  $\theta$  at  $t_c$  feels the tilt of potential and must decide to which of the two vacua  $\theta_{min} = 0$  or  $\theta_{min} = 2\pi$  it should roll down at the beginning of oscillation. We deal with the situation when at the beginning of inflation the Universe contains the phase  $\theta_{N_{max}} < \pi$  and hence the final state of the main part of the present particle horizon is  $\theta_{min} = 0$ . On the contrary, there will be the islands with  $\theta > \pi$ , which are the results of fluctuations (9) with positive sign. The phases inside that islands will move to the final state  $\theta_{min} = 2\pi$  after the triggering of the phase transition.

Thus, the strongly non-equilibrium distribution of the phase, causing by inflation, leads to the formation of islands with vacuum  $2\pi$  in the space with zero phase at the time moment  $t_c$ . Both states are separated by closed walls at  $\theta = \pi$ . The distribution on size for such closed domain walls resembles the size distribution of fluctuations (9) that consists on the steps with positive sign and move the phase to the region  $\theta > \pi$  during the inflational stage. The probability to have such fluctuations can be significant that provides us with the possibility to form quite large number of very big closed domain walls.

## 4 Results and Discussion

As we have seen in the preceding section all regions with phase  $\theta > \pi$  are converted into islands with vacuum  $\theta_{\min} = 2\pi$  surrounding by the closed walls. The size distribution of closed walls imprints the size distribution of domains filled with phase coming from fluctuations that have crossed the point  $\pi$  during the one dimensional brownian motion. The physical size that leaves the horizon during  $e$ -fold number  $N$  ( $N \leq N_{max}$ ) reads

$$l = H_i^{-1} e^N \quad (12)$$

This scale becomes comparable to the FRW particle horizon at the moment

$$t_h = H^{-1} e^{2N} \quad (13)$$

It is clear that we will start to observe the selfcollapse of a closed domain wall when its size is causally connected. Approximately at the same time the wall is acquiring spherical form due to the surface tension. Thus, if the amount of energy stored in a such vacuum configuration

$$E \approx \sigma t_h^2 \quad (14)$$

is large enough, the BH can be formed in the result of its selfcollapse. More definitely, the gravitation radius of configuration should exceed the minimal size up to which it can collapse. In our case the collapse of closed domain wall, coming from potential (1), changes on repulsion at the size comparable with wall's width<sup>17)</sup>. This establish cut off for the PBH's mass spectrum at the small masses range.

To evaluate numerically mass spectrum of PBHs we have to calculate the size distribution of domains that contains phases, which are at the range  $\theta > \pi$ . Suppose that at  $e$ -fold  $N = \ln(lH_i)$  before the end of inflation the volume  $V(\bar{\theta}, N)$  has been filled with phase value  $\bar{\theta}$ . Then at the  $e$ -fold  $N - 1$  the volume filled with average phase  $\bar{\theta}$  obeys following iterative expression<sup>21)</sup>

$$V(\bar{\theta}, N - 1) = e^3 V(\bar{\theta}, N) + (V_U(N) - e^3 V(\bar{\theta}, N)) P(\bar{\theta}, N - 1) \delta\theta, \quad (15)$$

here the  $V_U(N) \approx e^{3N} H_i^{-3}$  is the volume of the Universe at  $N$   $e$ -fold. We applied here distribution (10). Now one can easily calculate the size distribution of domains filled with appropriate value of phase corresponding to  $N$ , with the use of expression (15). For our numerical calculations we have chosen the following reasonable values for inflational Hubble constant  $H = 10^{13}$  GeV and for the radius of PNG potential  $f = 10^{14}$  GeV (see for example<sup>22)</sup>). Also we suppose that  $N_{max} = 60$  and that the total energy of the wall is smaller than the total energy of the medium inside it. The simulation has been performed for two cases that depend on the  $\Lambda$  and  $\theta_{N_{max}}$ . The results are following.

$\log_{10} \frac{M_{PBH}}{1g.}$	12	13	14	15	16	17	20	21	22
$\log_{10} n_{PBH}$	20.4	18.4	16.3	14.2	12.0	9.79	5.18	2.77	0.30

$$\Lambda = 10^8 \text{ GeV (see for example}^{22)}), \theta_{N_{max}} = 0.65$$

Table 1. The total number  $n_{PBH}$  of PBHs against their masses  $M_{PBH}$ .

The results of Table 1. are fitted in such a manner to satisfy the most stronger astrophysical constraints. It is known <sup>23)</sup> that only PBHs with masses larger then  $\approx 10^{15}$  g. can survive in respect to Hawking evaporation. The observations of diffused gamma ray background establish strong limit <sup>24)</sup> on the density fraction  $\Omega_{PBH} < 10^{-9}$  of PBHs with masses  $10^{14} \div 10^{15}$  g.. The next limit, which has to be checked is the limit on the abundance of PBHs with masses  $10^{12} \div 10^{13}$  g.. Although such light PBHs disappeared already due to Hawking radiation it could produce a large amount of entropy at the epoch of nucleosynthesis <sup>25)</sup>. Actually the constraints <sup>25)</sup> is valued for the PBH's mass range  $10^9 \div 10^{13}$  g. and reads  $\beta < 10^{-15}(10^9 g./M_{PBH})$ , here  $\beta$  is the density fraction of PBHs at the moment of their formation. It means that at the moment of full evaporation of such PBHs  $\tau = M_{PBH}^3/(g_* m_p^4)$  their contribution into the total density of the Universe should satisfy the following restriction  $\alpha < 10^{-4}$ . To check our spectrum we express the density fraction of PBHs at the moment of their full evaporation in the terms of  $n_{PBH}$ , it gives  $\alpha \simeq 10^{-8} n_{PBH} (1g./M_{PBH})^2$ . Thus in the case of Table 1. the entropy production limit <sup>25)</sup> is also satisfied. The more careful consideration of particle content of  $10^{10} \div 10^{13}$ g. PBH decay products shows that such PBHs are very effective sources of antiprotons at the period of their evaporation <sup>6)</sup>. According to the theory on non-equilibrium nucleosynthesis (see <sup>7)</sup> and references therein) it can change dramatically the abundance of such elements as  ${}^3He$ ,  ${}^6Li$  and  ${}^7Li$ . This fact establishes the most stringent constraints on the density fraction of PBHs in the mass range  $10^{10} \div 10^{13}$ g. <sup>7)</sup> <sup>3)</sup>. Thus, taking into account the effects of non-equilibrium nucleosynthesis of  ${}^3He$  and lithium we have got following limits on the density fraction of PBHs with mass range indicated above:  $\alpha_{3He} < 10^{-15} \Omega_b (M_{PBH}/1g.) k^{-1}$ ,  $\alpha_{Li} < 10^{-19} \Omega_b (M_{PBH}/1g.)^{5/4} k^{-1}$ , where the  $\Omega_b$  is the fraction of baryon density and  $k$  varies from 1/4 to 1/6 for PBHs of different masses <sup>7)</sup>. These limits are also satisfied our model in Table 1. As well we can see from the Table 1. that the number of BH with masses more than  $10^{22}$  g. is negligible, whereas BHs with masses smaller than  $10^{12}$  g. were not produced at all because of their small gravitational radius.

Another interesting case is small  $\Lambda$  limit that comes from QCD <sup>9)</sup>. It gives rise to massive BH.

$\log_{10} \frac{M_{PBH}}{1g.}$	28.3	29.5	30.6	31.8	33.0	34.1	35.3	36.4
$\log_{10} n_{PBH}$	17.4	13.0	11.1	9.2	7.2	5.1	3.0	0.66

$$\Lambda = 1\text{GeV (see for exapmle }^9), \theta_{N_{max}} = 0.8$$

Table 2. The total number  $n_{PBH}$  of PBHs against their masses  $M_{PBH}$ .

In this case the wall width is very large and so that only rather massive BHs could be formed. This fact can be applied for explanation of giant BHs origin in the centres of galaxies (see for example <sup>27)</sup>) The maximal PBH masses in the Table 2 are at least by 3–4 orders of the magnitude smaller than the masses of BHs, assumed to be present in the centres of galaxies. However the account for possible strong concentration of small mass PBHs around close-to-maximal mass PBHs can provide the evolution and collapse of PBH systems into black holes with the mass, exceeding by several orders of the magnitude the maximal PBH mass. It provides the possibility to evolve the approach to the origin of AGNs on the base of our model.

To complete our discussion let us mention following. We assumed that the coherent field oscillations around the true vacua, which start after the triggering of phase transition with the amplitude proportional to the residual between the local value of phase and the local vacuum phase, are damped due to relatively strong dissipation, so that no energy density is stored in the form of such oscillations. If the dissipation is not effective, as it is the case for invisible axion <sup>9)</sup>, the inhomogeneity of the phase

<sup>3)</sup>Here we take into account only the standard model set of particle spices. The invoking of the some supersymmetric extensions of the standard model with a large number of unstable moduli fields which could come from evaporating PBHs also would influence on the primordial chemical content of the Universe (see for example <sup>26)</sup>).

distribution before the phase transition results in the inhomogeneity of the energy density distribution of coherent field oscillations. Then such energy density contributes into the total cosmological density and its large scale inhomogeneity induces effects of anisotropy of relic radiation. In this case one can strongly constrain the model parameters from the observational upper limits on the total cosmological density and on the possible anisotropy of thermal background radiation induced by isocurvature perturbations<sup>28)</sup>.

Leaving a more completely study of the last two issues to future publications, we can conclude that the non-equilibrium dynamics of scalar field with effectively flat direction can give a new spin to the PBHs formation mechanisms applying collapse of domain walls.

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